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CMB anisotropy from baryogenesis by a scalar field

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Abstract

We study the cosmic microwave background (CMB) anisotropy in the scenario where the baryon asymmetry of the universe is generated from a condensation of a scalar field. In such a scenario, the scalar condensation may acquire fluctuation during the inflation which becomes a new source of the cosmic density perturbations. In particular, the primordial fluctuation of the scalar condensation may induce correlated mixture of the adiabatic and isocurvature fluctuations. If the scalar condensation decays before it completely dominates the universe, the CMB angular power spectrum may significantly deviate from the conventional adiabatic result. Such a deviation may be observed in the on-going MAP experiment.

1. Introduction

In modern particle cosmology, one of the most important issues is to understand the origin of the baryon asymmetry of the universe. In particular, assuming inflation [1] as a solution to the horizon, flatness, and other cosmological problems and as a seed for density fluctuations for later structure formation, we are obliged to adopt scenarios where the baryon asymmetry is generated after the reheating.

One of the main branches of baryogenesis is to use a primordial condensation of a scalar field as a source of the baryon-number asymmetry of the universe. In such a class of scenarios, there exists a primordial condensation of a scalar field and it decays at a later stage of the evolution of the universe. Then, baryon asymmetry is generated by the decay of the scalar field. If such a scalar field exists, one of the Sakharov's three conditions for baryogenesis, i.e., the out-of-equilibrium condition, is easily satisfied by assuming non-vanishing initial amplitude of the scalar field. In particular, in cosmological scenarios based on supersymmetric models, there exist various possible candidates of such a scalar field and many efforts have been made to study such scenarios.

Among various possibilities, probably one of the most well-motivated candidates of such scalar fields is the superpartner of the right-handed neutrino, the right-handed sneutrino. Assuming the seesaw mechanism [2] as a origin of very tiny neutrino masses which are suggested from the neutrino-oscillation experi-

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ments [3],[1] the right-handed sneutrinos inevitably exist in the supersymmetric models. Such a right-handed sneutrino can be a source of the baryon-number asymmetry [8,9]; if one of the right-handed neutrinos has non-vanishing amplitude in the early universe, its decay may generate lepton-number asymmetry which is converted to the baryon-number asymmetry via the sphaleron interaction [10]. (In fact, this is a supersymmetric extension of the leptogenesis scenario proposed by Fukugita and Yanagida [11].)

In addition, Affleck–Dine mechanism for baryogenesis [12] is another possibility. In the Affleck–Dine scenario, using the baryon-number violating operator as a source, non-vanishing baryon number is generated while a scalar partner of quarks, called Affleck–Dine field, is oscillating.

Since the baryogenesis by a scalar condensation is attractive and well-motivated, in particular, in supersymmetric models, it is important to consider how the scenario can be experimentally tested. As we will discuss below, one possible effect of such a scalar field is on the cosmological density perturbations; if a condensation of the scalar field in the early universe is the source of the present baryon asymmetry, some signal may be imprinted in the cosmic microwave background (CMB) since the significant fraction of the CMB radiation we observe today may also originate in the scalar-field condensation. In particular, the amplitude of the scalar field may acquire sizable fluctuation during the inflation and such a fluctuation affects the CMB anisotropy.

Thus, in this Letter, we study the cosmic density perturbations in the scenario of the baryogenesis by a scalar field. Although our conclusions are quite general to a large class of scenarios, we mostly concentrate on the scenario of the sneutrino-induced leptogenesis [8,9] to make our discussion clearer. In particular, we study effects of the primordial fluctuation of the scalar field on the CMB anisotropy. If the primordial fluctuation exists in the amplitude of the scalar field, correlated mixture of the adiabatic and isocurvature fluctuations may be generated which may significantly affect the CMB angular power spectrum. With the precise measurement of the CMB angular power spectrum by the MAP experiment, some signal of the scalar-field-induced leptogenesis may be observed.

[1] Even for Dirac neutrinos, an analogue of the seesaw mechanism is possible [4,5] and the leptogenesis is possible [6] via the decay of a scalar condensate [7].

2. The leptogenesis scenario

Let us start by introducing the scenario we consider. For our argument, there are two scalar fields which play important roles; one is the inflaton field χ and the other is the (lightest) right-handed sneutrino. In the following, we denote the *lightest* right-handed scalar neutrino as $\tilde{N}$. We assume the inflation so that the universe starts with a de Sitter epoch when the universe is dominated by the potential energy of the inflaton field χ. During the inflation, the right-handed sneutrino is assumed to have a non-vanishing amplitude $\tilde{N}_{\rm init}$. We treat $\tilde{N}_{\rm init}$ as a free parameter in the following discussion. After the inflation, the inflaton field starts to oscillate and then decays. The right-handed neutrino also starts to oscillate when the expansion rate of the universe becomes comparable to the Majorana mass M_N. (We assume that the expansion rate during the inflation is larger than M_N so that $\tilde{N}$ starts to oscillate after the inflation.) Then, the right-handed sneutrino decays when $H \sim \Gamma_N$, where Γ_N is the decay rate of $\tilde{N}$.

One important point in this scenario is that the CMB radiation we observe today has two origins; the inflaton field and the right-handed sneutrino. This is because the decays of the inflaton and the right-handed sneutrino both convert the energy densities stored in the scalar condensation into that of radiation. For our study, it is convenient to distinguish the radiation originating in the inflaton field from that originating in the right-handed sneutrino. We denote them as γ_χ and $\gamma_{\tilde{N}}$, respectively.[2] Energy fraction of the radiation from $\tilde{N}$ depends on the initial amplitude $\tilde{N}_{\rm init}$ as well

[2] In fact, we cannot neglect the momentum transfer between these photons and they cannot be defined separately. In particular, velocity perturbations of these photons should be the same since the mean free path of the photon is much shorter than the horizon scale. If we only discuss the behavior of δ_{γ_χ} for the super-horizon modes, however, the following discussions are valid since the velocity perturbations are suppressed by the factor $k\tau$ relative to the density perturbations. In a rigorous sense, γ_χ and $\gamma_{\tilde{N}}$ should be understood as representatives of the components produced from the decay products of the inflaton field and $\tilde{N}$, respectively.

as the decay rate of the inflaton χ (denoted as Γ_χ), decay rate of $\tilde{N}$ and the mass of the right-handed (s)neutrino.

If $\Gamma_\chi \lesssim \Gamma_N$, $\tilde{N}$ decays when the universe is dominated by the inflaton condensation. In particular, when $\Gamma_N \lesssim H \lesssim M_N$, energy densities of $\tilde{N}$ and χ are related as[3]

$$\rho_{\tilde{N}}/\rho_\chi \sim \tilde{N}_{\rm init}^2/M_*^2, \qquad \Gamma_N \lesssim H \lesssim M_N. \tag{2.1}$$

(Here and hereafter, ρ_X denotes the energy density of the component X.) Using this relation, we can evaluate the energy density of the right-handed sneutrino at the time of its decay; when $H \sim \Gamma_N$, $\rho_\chi \sim \Gamma_N^2 M_*^2$ and hence $\rho_{\tilde{N}} \sim \Gamma_N^2 \tilde{N}_{\rm init}^2$. After the decay, energy density stored in the sneutrino condensation is converted to that of radiation, so the universe is filled with the radiation (denoted as γ_N) as well as the inflaton oscillation. Using the relations $\rho_\chi \propto a^{-3}$ and $\rho_{\tilde{N}} \propto a^{-4}$, we obtain the energy densities of each components at $H \sim \Gamma_\chi$ as $\rho_\chi \sim \Gamma_\chi^2 M_*^2$ and $\rho_{\gamma_{\tilde{N}}} \sim \rho_\chi (\Gamma_\chi/\Gamma_N)^{2/3} (\tilde{N}_{\rm init}/M_*)^2$. Thus, after the decay of the inflaton field, the energy fraction of the radiation from the right-handed sneutrino is estimated as

$$f_{\gamma_{\tilde{N}}} \sim \frac{(\Gamma_\chi/\Gamma_N)^{2/3} (\tilde{N}_{\rm init}/M_*)^2}{1 + (\Gamma_\chi/\Gamma_N)^{2/3} (\tilde{N}_{\rm init}/M_*)^2}, \qquad \Gamma_\chi \lesssim \Gamma_N. \tag{2.2}$$

Here, f_{γ_X} is the energy fraction of γ_X after the decays of the inflaton and $\tilde{N}$, and $f_{\gamma_\chi} + f_{\gamma_{\tilde{N}}} = 1$. (Here and hereafter, $\gamma_X = \gamma_{\tilde{N}}$ or γ_χ.)

If $\Gamma_\chi \gtrsim \Gamma_N$, on the contrary, $\tilde{N}$ decays after the inflaton decay. In this case, it is convenient to define the following quantity:

$$\tilde{N}_{\rm eq} \sim \begin{cases} (\Gamma_N/M_N)^{1/4} M_*, & M_N < \Gamma_\chi, \\ (\Gamma_N/\Gamma_\chi)^{1/4} M_*, & M_N > \Gamma_\chi. \end{cases} \tag{2.3}$$

[3] When the initial amplitude of $\tilde{N}$ is larger than M_*, inflation is caused by the energy density of $\tilde{N}$. If this happens, all the components generated from χ is washed out and $f_{\gamma_{\tilde{N}}} \to 1$. If the e-folding number of the secondary inflation caused by $\tilde{N}$ is small enough, the following discussions are unchanged since the density fluctuation for the scale corresponding to the CMB anisotropy we will discuss is generated from the primary inflation caused by χ. If the initial amplitude is much larger than M_*, only the energy density from the decay of the right-handed neutrino is relevant and the perturbation becomes purely adiabatic [13].

If $\tilde{N}_{\rm init} \sim \tilde{N}_{\rm eq}$, $\rho_{\gamma_\chi} \sim \rho_{\tilde{N}}$ is realized when $H \sim \Gamma_{\tilde{N}}$. Thus, when $\tilde{N}_{\rm init} \lesssim \tilde{N}_{\rm eq}$, $\tilde{N}$ decays in the γ_χ-dominated universe and hence

$$f_{\gamma_{\tilde{N}}} \sim \frac{(\tilde{N}_{\rm init}/\tilde{N}_{\rm eq})^2}{1 + (\tilde{N}_{\rm init}/\tilde{N}_{\rm eq})^2}, \qquad \Gamma_N \lesssim \Gamma_\chi,\ \tilde{N}_{\rm init} \lesssim \tilde{N}_{\rm eq}. \tag{2.4}$$

On the contrary, if $\tilde{N}_{\rm init} \gtrsim \tilde{N}_{\rm eq}$, the right-handed sneutrino decays after it dominates the universe and we obtain

$$f_{\gamma_{\tilde{N}}} \sim \frac{(\tilde{N}_{\rm init}/\tilde{N}_{\rm eq})^{8/3}}{1 + (\tilde{N}_{\rm init}/\tilde{N}_{\rm eq})^{8/3}}, \qquad \Gamma_N \lesssim \Gamma_\chi,\ \tilde{N}_{\rm eq} \lesssim \tilde{N}_{\rm init}. \tag{2.5}$$

When $\tilde{N}$ decays, lepton-number asymmetry is also generated. Such a lepton-number asymmetry is converted to the baryon-number asymmetry due to the sphaleron process. The resultant baryon-number asymmetry depends on whether $\tilde{N}$ decays before or after the $\tilde{N}$-dominated universe is realized. If $\tilde{N}$ decays later than the inflaton decay, the baryon-to-entropy ratio is given by [9]

$$\frac{n_B}{s} \simeq 0.24 \times 10^{-10} f_{\gamma_{\tilde{N}}} \delta_{\rm eff} \left(\frac{T_N}{10^6\ {\rm GeV}}\right) \left(\frac{m_{\nu_3}}{0.05\ {\rm eV}}\right), \qquad \Gamma_N < \Gamma_\chi, \tag{2.6}$$

where T_N is the temperature at the epoch of the decay of $\tilde{N}$, m_{ν_3} the mass of the heaviest (left-handed) neutrino mass. (So, if $\tilde{N}$ decays after dominating the universe, T_N becomes the reheat temperature due to the decay of $\tilde{N}$.) In addition, in the basis where the Majorana mass matrix for the right-handed neutrinos $\widehat{M}$ is real and diagonalized, the effective CP violating phase is given by

$$\delta_{\rm eff} = \frac{\langle H_u \rangle^2}{m_{\nu_3}} \frac{{\rm Im}[\hat{h}\hat{h}^\dagger \widehat{M}^{-1} \hat{h}^* \hat{h}^T]_{11}}{[\hat{h}\hat{h}^\dagger]_{11}}, \tag{2.7}$$

where $\hat{h}_{i\alpha}$ is the neutrino Yukawa matrix with i and α being the generation indices for the right-handed and left-handed neutrinos, respectively. Notice that, with a maximum CP violation, $\delta_{\rm eff} \sim 1$. On the contrary, if $\Gamma_\chi < \Gamma_N$, $\tilde{N}$ decays before the decay of the inflaton field. In this case, radiation-dominated universe is realized much later than the epoch of the sneutrino decay. In this case, the baryon-to-entropy ratio is given by

$$\frac{n_B}{s} \simeq 0.24 \times 10^{-10} \times \delta_{\rm eff}\left(\frac{T_{\rm R}}{10^6\ {\rm GeV}}\right)\left(\frac{m_{\nu_3}}{0.05\ {\rm eV}}\right)\left(\frac{\widetilde{N}_{\rm init}}{M_*}\right)^2, \quad \Gamma_\chi < \Gamma_N, \tag{2.8}$$

where $T_{\rm R}$ is the reheating temperature due to the decay of the inflaton field.[4]

As one can see, large enough baryon asymmetry ($n_B/s \sim 8 \times 10^{-11}$ [15]) can be generated in this scenario with relatively low reheating temperature. In particular, even if we impose the upper bound on the reheating temperature to avoid the overproduction of the gravitinos, it is possible to generate large enough baryon asymmetry. For $\Gamma_N < \Gamma_\chi$, requiring $T_N \lesssim 10^{9-10}$ GeV in order not to overproduce gravitinos [16], $n_B/s \sim 10^{-10}$ can be realized when $f_{\gamma_{\widetilde{N}}} \gtrsim 10^{-3} - 10^{-2}$. For $\Gamma_\chi < \Gamma_N$, much smaller $f_{\gamma_{\widetilde{N}}}$ is possible to generate enough baryon asymmetry even if we require $T_{\rm R} \lesssim 10^{9-10}$ GeV to avoid the gravitino problem. If the above gravitino constraint can be somehow neglected, T_N (or $T_{\rm R}$) can be higher and $n_B/s \sim 10^{-10}$ can be realized even with smaller value of $f_{\gamma_{\widetilde{N}}}$.[5]

3. Density perturbations

As we have seen, in the scenario we consider, baryon asymmetry as well as some fraction of the CMB radiation are generated from the decay product of $\widetilde{N}$. Thus, if $\widetilde{N}$ has a primordial fluctuation, such a fluctuation also becomes the source of the cosmic density perturbations. Most importantly, such a fluctuation affects the CMB anisotropy which is now being measured very precisely with various experiments [17–21].

Fluctuation of $\widetilde{N}$ is primarily induced during the inflation; assuming that the (effective) mass of the right-handed sneutrino during the inflation is smaller than the expansion rate during the inflation $H_{\rm inf}$, the primordial fluctuation of $\widetilde{N}$ is estimated as[6]

$$\delta\widetilde{N}_{\rm init} = \frac{H_{\rm inf}}{2\pi}. \tag{3.1}$$

If $\delta\widetilde{N}_{\rm init}$ is non-vanishing, this becomes a source of the entropy between components generated from the decay products of χ and $\widetilde{N}$. It is convenient to define the parameter

$$S_{\widetilde{N}\chi}^{(\delta\widetilde{N})} \equiv \delta_{\widetilde{N}}^{(\delta\widetilde{N})} - \delta_{\chi}^{(\delta\widetilde{N})}, \tag{3.2}$$

where $\delta_X \equiv \delta\rho_X/\rho_X$ with ρ_X being the energy density of the component X and $\delta\rho_X$ its perturbation in the Newtonian gauge. Here, the right-hand side is evaluated after $\widetilde{N}$ and χ both start to oscillate, and the superscript "$(\delta\widetilde{N})$" is for variables generated from the primordial fluctuation of $\widetilde{N}$. Solving the equations of motions for the scalar fields, we obtain [22,23]

$$S_{\widetilde{N}\chi}^{(\delta\widetilde{N})} = \frac{2\delta\widetilde{N}_{\rm init}}{\widetilde{N}_{\rm init}}, \tag{3.3}$$

where we assumed that the potential of the scalar fields are dominated by the quadratic terms. After the decays of $\widetilde{N}$ and χ, $S_{\widetilde{N}\chi}^{(\delta\widetilde{N})}$ becomes entropy between components generated from the decay products of $\widetilde{N}$ and χ for the perturbations with wavelength much longer than the horizon scale.

With the non-vanishing value of the primordial fluctuation of $\widetilde{N}$ given in Eq. (3.1), it is important to note that there are two independent sources of the perturbations in our case; fluctuation of the inflaton field and that of the right-handed sneutrino. In the framework of the linear perturbation theory, we can discuss effects of these perturbations separately. In addition, in discussing the CMB angular power spectrum C_l, effects of the two perturbations can be treated separately since we assume no correlation between these two fields. Thus, the total angular power spectrum is

[4] Eqs. (2.6) and (2.8) are applicable even if the parametric-resonant decay of $\widetilde{N}$ is effective at some stage. If the initial amplitude of $\widetilde{N}$ is larger than $\widetilde{N}_{\rm PR} \sim (10\text{–}100) M_N/h$, the amplitude decreases down to $\widetilde{N}_{\rm PR}$ due to parametric resonance just after $\widetilde{N}$ starts to oscillate, where h is the neutrino Yukawa coupling for $\widetilde{N}$ [14]. In this case, $\widetilde{N}_{\rm init}$ in Eq. (2.8) (and in other formulae) should be replaced by $\widetilde{N}_{\rm PR}$.

[5] When T_N (or $T_{\rm R}$) becomes high, thermal effects on the potential of $\widetilde{N}$ may become significant and may change the discussion given in this section. The thermal effects may be, however, neglected if M_N, T_N (or $T_{\rm R}$) larger than is assumed.

[6] During the inflation, the effective mass of $\widetilde{N}$ may become as large as the expansion rate of the universe. If so, the primordial fluctuation of $\widetilde{N}$ is extremely suppressed. In some class of models, however, this is not the case, like in the no-scale type models and D-term inflation model [24].

given in the form

$$C_l = C_l^{(\delta\chi)} + C_l^{(\delta\widetilde{N})}, \tag{3.4}$$

where $C_l^{(\delta\chi)}$ and $C_l^{(\delta\widetilde{N})}$ are the contributions from the primordial fluctuations in the inflaton and $\widetilde{N}$, respectively. The inflaton contribution $C_l^{(\delta\chi)}$ is known to become the adiabatic result (with relevant spectral power index). Thus, let us consider the second term $C_l^{(\delta\widetilde{N})}$.

In discussing the cosmic density perturbations induced from $\delta\widetilde{N}_{\rm init}$, we treat the radiations originated from these fields separately; γ_χ for the radiation from the inflaton and $\gamma_{\widetilde{N}}$ which is that from the right-handed sneutrino. Density perturbations for these components are defined separately: $\delta_{\gamma X} \equiv \delta\rho_{\gamma X}/\rho_{\gamma X}$ and $V_{\gamma X}$ are density and velocity fluctuations of γ_X with X being χ and $\widetilde{N}$. (Hereafter, we use the Newtonian gauge. We follow the notation and convention of [25].)

In the following, we follow the evolution of the density perturbations of various components solving the relevant Einstein and Boltzmann equations. In particular, we are interested in evolutions of the density perturbations when $H \gtrsim \Gamma_N$ and properties of the density perturbations just after the decay of $\widetilde{N}$. In this case, we consider the evolution of the perturbations in the universe with a very high temperature where various charged particles are thermally produced. Then, the radiation component becomes locally isotropic and hence the anisotropic stress perturbation can be neglected.

Denoting the perturbed line element as

$$ds^2 = a^2\big[-(1+2\Psi)\,d\tau^2 + (1+2\Phi)\delta_{ij}\,dx^i\,dx^j\big], \tag{3.5}$$

with τ being the conformal time and a being the scale factor, the Poisson equation for the metric perturbation is given by

$$k^2\Phi = 4\pi G a^2 \rho_{\rm T}\left[\delta_{\rm T} + 3\frac{a'}{a}(1+w_{\rm T})V_{\rm T}/k\right], \tag{3.6}$$

where k is the comoving momentum, G the Newton constant, and the "prime" denotes the derivative with respect to the conformal time τ. In addition, the subscript "T" is for the total matter and $w_{\rm T}$ denotes the equation-of-state parameter of the total matter. On the other hand, neglecting the contribution from the anisotropic stress tensor, evolution of the density and velocity perturbations in the Fourier space are given by [25]

$$\delta'_{\gamma X} = -\frac{4}{3}kV_{\gamma X} - 4\Phi',$$
$$V'_{\gamma X} = k\left(\frac{1}{4}\delta_{\gamma X} + \Psi\right). \tag{3.7}$$

In addition, in the situation we consider, the relation $\Phi = -\Psi$ holds since the anisotropic stress tensor is small enough. In the following, we use this relation to eliminate Φ.

Using Eqs. (3.6) and (3.7), we can discuss the evolution of the cosmic density perturbations induced from the primordial fluctuation of the amplitude of the right-handed sneutrino condensation. In this case, we neglect the adiabatic fluctuations generated by the inflaton fluctuation. In addition, since we are interested in the perturbations whose wavelength is much longer than the horizon scale (at the epoch of the baryogenesis and reheating), we can expand the solution as a function of $k\tau$.

Evolution of δ_{γ_χ} is quite simple. Just after the inflation, the right-handed sneutrino is the sub-dominant component and hence $\Psi^{(\delta\widetilde{N})}$ vanishes as $\tau \to 0$. In addition, since there is no primordial perturbation in the inflaton sector for this mode, $\delta_{\gamma_\chi}^{(\delta\widetilde{N})} \to 0$ as $\tau \to 0$. Thus, we obtain

$$\delta_{\gamma_\chi}^{(\delta\widetilde{N})} = 4\Psi^{(\delta\widetilde{N})} + O\big(k^2\tau^2\big). \tag{3.8}$$

This relation holds at any moment of the evolution of the universe.

Using the above relation, the entropy between the radiation and the baryon components can be evaluated. We calculate the entropy in the radiation-dominated universe after the decay of both χ and $\widetilde{N}$. In the radiation-dominated universe [25]

$$\Delta_{\rm T} = O\big(k^2\tau^2\big), \qquad \delta_{\rm T} = -2\Psi_{\rm RD},$$
$$V_{\rm T} = \frac{1}{2}\Psi_{\rm RD}k\tau. \tag{3.9}$$

(Here and hereafter, we only denote the leading contribution to the perturbations.) Then, using the relation

$$\Delta_{\rm T} = f_{\gamma_\chi}\Delta_{\gamma_\chi} + f_{\gamma_{\widetilde{N}}}\Delta_{\gamma_{\widetilde{N}}}, \tag{3.10}$$

with $\Delta_{\gamma X} = \delta_{\gamma X} + 4(a'/a)V_{\rm T}/k$, we obtain

$$\Delta^{(\delta\tilde{N})}_{\gamma_{\tilde{N}}} = -\frac{f_{\gamma_\chi}}{f_{\gamma_{\tilde{N}}}}\Delta^{(\delta\tilde{N})}_{\gamma_\chi} = -6\frac{f_{\gamma_\chi}}{f_{\gamma_{\tilde{N}}}}\Psi^{(\delta\tilde{N})}_{\rm RD}. \tag{3.11}$$

Using the fact that the entropy between any component produced from $\tilde{N}$ and that from the inflaton field is conserved, we can relate $\Psi^{(\delta\tilde{N})}_{\rm RD}$ with $S^{(\delta\tilde{N})}_{\tilde{N}\chi}$; with relation $S^{(\delta\tilde{N})}_{\tilde{N}\chi} = \frac{3}{4}(\Delta^{(\delta\tilde{N})}_{\gamma_{\tilde{N}}} - \Delta^{(\delta\tilde{N})}_{\gamma_\chi})$, we obtain

$$\Psi^{(\delta\tilde{N})}_{\rm RD} = -\frac{2}{9}f_{\gamma_{\tilde{N}}}S^{(\delta\tilde{N})}_{\tilde{N}\chi} = -\frac{4}{9}f_{\gamma_{\tilde{N}}}\frac{\delta\tilde{N}_{\rm init}}{\tilde{N}_{\rm init}}. \tag{3.12}$$

Thus, if $\tilde{N}$ decays after it dominates the universe, $f_{\gamma_{\tilde{N}}} \simeq 1$ and hence the metric perturbation becomes comparable to the primordial entropy perturbation. On the other hand, if $f_{\gamma_{\tilde{N}}} \ll 1$, the metric perturbation becomes negligibly small.

Since the baryon asymmetry is generated from $\tilde{N}$, the density fluctuation in the baryonic component is given by

$$\Delta^{(\delta\tilde{N})}_{b} = \frac{3}{4}\Delta^{(\delta\tilde{N})}_{\gamma_{\tilde{N}}} = -\frac{9}{2}\frac{f_{\gamma_\chi}}{f_{\gamma_{\tilde{N}}}}\Psi^{(\delta\tilde{N})}_{\rm RD}. \tag{3.13}$$

Thus, the entropy between the baryon and the radiation is given by

$$\begin{aligned} S^{(\delta\tilde{N})}_{b\gamma} &= \Delta^{(\delta\tilde{N})}_{b} - \frac{3}{4}\Delta^{(\delta\tilde{N})}_{\rm T} \\ &= -\frac{9}{2}\frac{f_{\gamma_\chi}}{f_{\gamma_{\tilde{N}}}}\Psi^{(\delta\tilde{N})}_{\rm RD} = -\frac{9(1-f_{\gamma_{\tilde{N}}})}{2f_{\gamma_{\tilde{N}}}}\Psi^{(\delta\tilde{N})}_{\rm RD}. \end{aligned} \tag{3.14}$$

Thus, the correlated mixture of the adiabatic and isocurvature fluctuations is generated. In addition, if the right-handed sneutrino once dominates the universe, $f_{\gamma_{\tilde{N}}} \to 1$ and hence the perturbation becomes adiabatic [22,23,26,27]. (Thus, in this case, $\tilde{N}$ may play the role of the so-called "curvaton" field.) On the contrary, if the inflaton field and its decay products always dominate the universe, $f_{\gamma_{\tilde{N}}} \to 0$ and the perturbation becomes isocurvature. Thus, the interesting spectrum would be obtained when the inflaton and the right-handed sneutrino both produce significant amount of the resultant radiation.

4. CMB angular power spectrum

Now we discuss the CMB angular power spectrum induced from the primordial fluctuation of $\tilde{N}$. For this purpose, we parameterize the entropy perturbation induced from $\delta\tilde{N}_{\rm init}$ as

$$S^{(\delta\tilde{N})}_{b\gamma} = \kappa_b\Psi^{(\delta\tilde{N})}_{\rm RD}. \tag{4.1}$$

Then, using Eqs. (2.2), (2.4), (2.5) and (3.14), the κ_b parameter is estimated as

$$\kappa_b \sim -\frac{9}{2}\times\begin{cases} (\Gamma_\chi/\Gamma_N)^{-2/3}\big(\tilde{N}_{\rm init}/M_*\big)^{-2}, \\ \quad \Gamma_\chi \lesssim \Gamma_N, \\ \big(\tilde{N}_{\rm init}/\tilde{N}_{\rm eq}\big)^{-2}, \\ \quad \Gamma_N \lesssim \Gamma_\chi,\ \tilde{N}_{\rm init} \lesssim \tilde{N}_{\rm eq}, \\ \big(\tilde{N}_{\rm init}/\tilde{N}_{\rm eq}\big)^{-8/3}, \\ \quad \Gamma_N \lesssim \Gamma_\chi,\ \tilde{N}_{\rm eq} \lesssim \tilde{N}_{\rm init}, \end{cases} \tag{4.2}$$

Thus, in our case, the κ_b parameter varies from $-\infty$ to 0. If $\tilde{N}_{\rm init} \ll \tilde{N}_{\rm eq}$, the κ_b parameter goes to $-\infty$ and hence the density perturbation becomes (baryonic) isocurvature. On the contrary, when $\tilde{N}_{\rm init} \gg \tilde{N}_{\rm eq}$, $\kappa_b \sim 0$ and the adiabatic density perturbation is obtained from the primordial fluctuation of $\tilde{N}$. The most interesting case is $\tilde{N}_{\rm init} \sim \tilde{N}_{\rm eq}$, where $\kappa_b \sim O(0.1\text{–}1)$. In this case, correlation between adiabatic and isocurvature perturbations becomes the most effective. For the case where $\Gamma_N < \Gamma_\chi$, the requirement of enough baryon asymmetry with the gravitino constraint give $f_{\gamma_{\tilde{N}}} \gtrsim 10^{-3} - 10^{-2}$, and hence $\kappa_b \gtrsim -1000$. If $\Gamma_\chi < \Gamma_N$, smaller κ_b is possible. (See Eqs. (2.6) and (2.8).)

In Fig. 1, we plot the CMB angular power spectrum induced from the primordial fluctuation of the right-handed sneutrino $C_l^{(\delta\tilde{N})}$. (Here and hereafter, we assume there is no entropy perturbation in the cold dark matter (CDM) sector.[7]) Notice that the lines with $\kappa_b = 0$ and $\kappa_b = -\infty$ coincide with the results with purely adiabatic and isocurvature density perturbations, respectively. For a general case, however, the angular power spectrum has a unique structure. In particular, as one can see, the acoustic peaks are suppressed relative to the Sachs–Wolfe (SW) tail as $|\kappa_b|$ increases.

As we discussed, the total angular power spectrum is given by the sum of the inflaton and the $\tilde{N}$ contributions, as shown in Eq. (3.4). Here, $C_l^{(\delta\chi)}$ is the inflaton contribution which is parameterized by

[7] If the lightest superparticle (LSP) is the CDM, this is the case since T_N (or $T_{\rm R}$) is much higher than the freeze-out temperature of the LSP. If the origin of the CDM is not the same as that of the CMB radiation, correlated entropy perturbation may also arise in the CDM sector. For such a case, see, for example, Ref. [23].

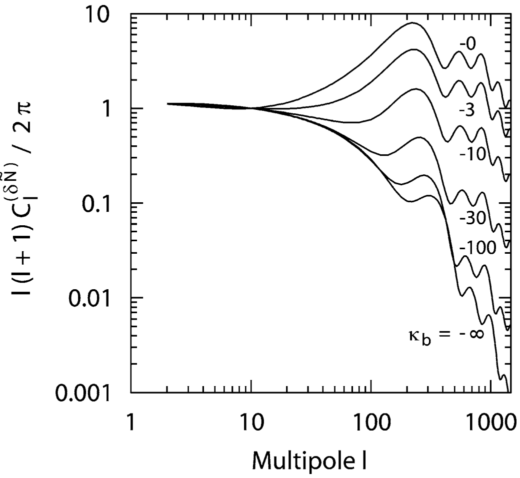

Fig. 1. The CMB angular power spectrum from the primordial fluctuation of the right-handed sneutrino. The κ_b parameter is taken as 0, -3, -10, -30, -100 and $-\infty$ from above. We consider the flat universe with $\Omega_b h^2 = 0.02$, $\Omega_m = 0.3$, and $h = 0.65$ [15] and the initial power spectral indices for primordial density perturbations are all assumed to be 1 (i.e., we adopt scale-invariant initial power spectra). The overall normalizations are taken as $[l(l+1)C_l^{(\delta\tilde{N})}/2\pi]_{l=10} = 1$.

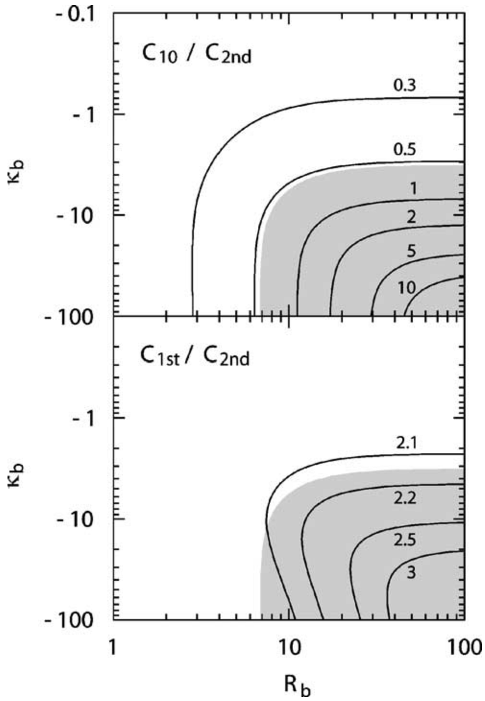

Fig. 2. The ratios $C_{10}/C_{2\mathrm{nd}}$ (top) and $C_{1\mathrm{st}}/C_{2\mathrm{nd}}$ (bottom) on the R_b vs. κ_b plane. (The numbers in the figure are the corresponding ratios.) The shaded region corresponds to the parameter space with $\chi^2 > 84$. The cosmological parameters are the same as those used in Fig. 1.

the primordial metric perturbation generated from the inflaton perturbation $\Psi^{(\delta\chi)}$, which is given by [28]

$$\Psi_{\mathrm{RD}}^{(\delta\chi)} = \frac{4}{9}\left[\frac{H_{\mathrm{inf}}}{2\pi}\frac{3H_{\mathrm{inf}}^2}{V'_{\mathrm{inf}}}\right]_{k=aH_{\mathrm{inf}}}, \tag{4.3}$$

where V_{inf} is the inflaton potential and $V'_{\mathrm{inf}} \equiv (\partial V_{\mathrm{inf}}/\partial\chi)$, and the superscript $(\delta\chi)$ means that the corresponding variable is generated from the inflaton fluctuation. On the contrary, $C_l^{(\delta\tilde{N})}$ is from the primordial fluctuation of $\tilde{N}$, and is parameterized by $S_{\tilde{N}\chi}^{(\delta\tilde{N})}$ given in Eq. (3.2). Thus, the total angular power spectrum crucially depends on three parameters, κ_b, $\Psi^{(\delta\chi)}$, and $S_{\tilde{N}\chi}^{(\delta\tilde{N})}$. To parameterize the relative size between $\Psi^{(\delta\chi)}$ and $S_{\tilde{N}\chi}^{(\delta\tilde{N})}$, we define

$$R_b = S_{\tilde{N}\chi}^{(\delta\tilde{N})}/\Psi_{\mathrm{RD}}^{(\delta\chi)}. \tag{4.4}$$

(Notice that, with the above definition, $\Psi_{\mathrm{RD}}^{(\delta\tilde{N})} = -\frac{2}{9} f_{\gamma\tilde{N}} R_b \Psi_{\mathrm{RD}}^{(\delta\chi)}$.) The R_b-parameter depends on the initial amplitude of the right-handed sneutrino as well as the model of inflation. For the chaotic inflation model, for example, $R_b \simeq 0.6 M_*/\tilde{N}_{\mathrm{init}}$ [23].

Then, the shape of the total angular power spectrum is determined once the parameters R_b and κ_b (as well as other cosmological parameters) are fixed. As these parameters vary, the shape of the angular power spectrum changes as follows. As the R_b-parameter is increased, $C_l^{(\delta\tilde{N})}$ is more enhanced and hence the height of the acoustic peaks are suppressed relative to the SW tail. In addition, as $|\kappa_b|$ is increased, C_l at high multipole is suppressed relative to that at low one since the acoustic peaks of $C_l^{(\delta\tilde{N})}$ is suppressed in this case. In Fig. 2, we plot the ratios $C_{10}/C_{2\mathrm{nd}}$ and $C_{1\mathrm{st}}/C_{2\mathrm{nd}}$ on the R_b vs. κ_b plane.

As one can read off from Fig. 2, the shape of the angular power spectrum deviates from that from purely adiabatic density perturbations if R_b and κ_b are non-vanishing. Since the current data of the CMB angular power spectrum is well consistent with C_l calculated from the purely adiabatic density perturbations, C_l in our scenario becomes inconsistent with the experiments if the R_b and $|\kappa_b|$ become too large. In order to derive the constraint on these parameters, we calculate the goodness-of-fit parameter $\chi^2 = -2\ln L$, where L is the likelihood function, as a function of R_b and κ_b. In our calculation, the offset log-normal approximation is used [29], and we use a data set consisting of 65 data points; 24 from COBE [17], 19 from BOOMERanG [18], 13 from MAXIMA [19], and 9 from DASI [20]. Then, in Fig. 2, we shaded the region where $\chi > 84$, which corresponds to 95% C.L. excluded region for the χ^2 statistics with 64 degrees of freedom. As expected, the angular power spectrum at lower multipoles is enhanced relative to that at higher multipoles. In particular, the ratio $C_{10}/C_{2\mathrm{nd}} \simeq 0.25$ in the purely adiabatic case, and hence the height of the SW tail can be enhanced by factor ~ 2 in the case of the sneutrino leptogenesis. In addition, $C_{1\mathrm{st}}/C_{2\mathrm{nd}} \simeq 2.0$ in the purely adiabatic case, and hence the enhancement of the ratio $C_{1\mathrm{st}}/C_{2\mathrm{nd}}$ is 5% or so in this case.

Precise determination of the CMB angular power spectrum by the MAP experiment will provide stronger constraints on our scenario. At the MAP experiment, uncertainty in C_l for multipoles $l \lesssim 1000$ is expected to be dominated by the cosmic variance, and in this case the error of C_l is given by [30]

$$\delta C_l = \sqrt{\frac{2}{2l+1}}\, C_l. \tag{4.5}$$

Thus, error of single C_l may not be small. However, combining the informations derived from C_l with different l, uncertainties can be reduced. For example, if we use the data for $2 \leqslant l \leqslant 50$, $201 \leqslant l \leqslant 250$, and $526 \leqslant l \leqslant 575$ to estimate the heights of the SW tail (represented by C_{10}) and the first and acoustic peaks, the errors are estimated as $\delta C_{10}/C_{10} \simeq 1.5\%$, $\delta C_{1\mathrm{st}}/C_{1\mathrm{st}} \simeq 0.9\%$, and $\delta C_{2\mathrm{nd}}/C_{2\mathrm{nd}} \simeq 0.6\%$. With this accuracy, effects of isocurvature fluctuations will be observed if the isocurvature contribution enhances the height of the SW tail by a few % or so.

One important question is whether we can distinguish the case with correlated mixture of the adiabatic and isocurvature fluctuations from the case with uncorrelated isocurvature fluctuation (i.e., $\kappa_b \to \infty$). Since the angular power spectrum at low multipoles are suppressed relative to C_l at higher multipoles in both cases, two cases are indistinguishable only by determining the ratio $C_{10}/C_{2\mathrm{nd}}$ or $C_{1\mathrm{st}}/C_{2\mathrm{nd}}$. As seen in Fig. 2, however, contours of constant $C_{10}/C_{2\mathrm{nd}}$ and those of $C_{1\mathrm{st}}/C_{2\mathrm{nd}}$ on the R_b vs. κ_b plane are not parallel and hence these two cases can be, in principle, distinguished by simultaneously determining the ratios $C_{10}/C_{2\mathrm{nd}}$ and $C_{1\mathrm{st}}/C_{2\mathrm{nd}}$.

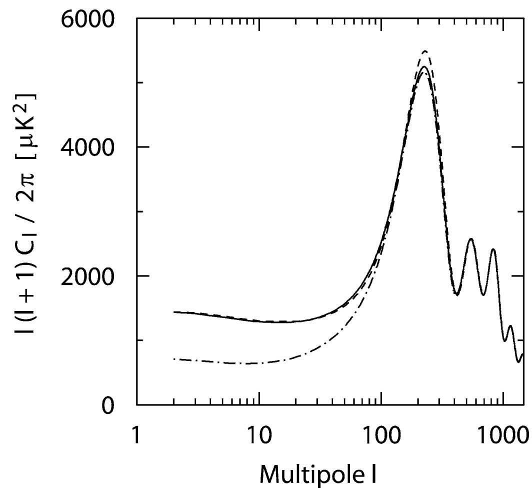

Fig. 3. The total CMB angular power spectrum with $(R_b, \kappa_b) = (100, -3)$ (dashed) and $(6.4, -100)$ (solid). The cosmological parameters are the same as those used in Fig. 1, and the normalizations are arbitrary. For comparison, we also plot the result in the purely adiabatic case (dot-dashed).

To discuss this issue, in Fig. 3, we plot the CMB angular power spectrum for $(R_b, \kappa_b) = (100, -3)$ (i.e., for the case with the correlated mixture of the adiabatic and isocurvature fluctuations) and $(6.4, -100)$ (i.e., for the case with uncorrelated isocurvature fluctuations). Here, we choose parameters such that the ratio $C_{10}/C_{2\mathrm{nd}}$ becomes the same for two cases. Even if the ratio $C_{10}/C_{2\mathrm{nd}}$ does not differ, the heights of the first acoustic peak are different; in this case, the ratio $C_{1\mathrm{st}}/C_{2\mathrm{nd}}$ differs about 5%, which is within the reach of the MAP experiment if the error in C_l mentioned before is realized. Thus, if the effect of the isocurvature mode is relatively large, an evidence of the correlated mixture of the adiabatic and isocurvature fluctuations may be observed by the MAP experiment.

5. Summary

In this Letter, we discussed the CMB anisotropy in the scenario where the baryon asymmetry of the universe originates in a scalar field condensation. We have seen that, in such a scenario, correlated mixture of the adiabatic and isocurvature fluctuations may be generated in particular when the decay product of the inflaton field and that of the scalar field (i.e., $\tilde{N}$ in our example) both significantly contribute to the present CMB radiation. With such a correlated mixture of fluctuations, the CMB angular power spectrum may be significantly affected and the on-going MAP experiment may observe the signal of the baryogenesis from the scalar-field condensation.

In our discussion, we used the sneutrino-induced leptogenesis as an example to make our discussion clearer. However, our discussion can be applied to a wide class of scenarios where the baryon asymmetry of the universe originates in a scalar-field condensation, like the Affleck–Dine scenario.[8]

Acknowledgement

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[8] The Affleck–Dine field is a complex scalar field and hence there may exist two independent fluctuations. Among various possibilities, one of the most interesting case is that the Affleck–Dine field significantly contributes to the CMB radiation today; if so, correlated mixture of the adiabatic and isocurvature fluctuation may arise. In this case, the oscillation of the Affleck–Dine field should be almost in the radial direction; otherwise, the resultant baryon-to-entropy ratio becomes too large. Then, the fluctuation in the radial direction provides the correlated mixture of the adiabatic and isocurvature fluctuations as in the $\tilde{N}$ case. The fluctuation in the phase direction provides purely isocurvature fluctuation in the baryonic sector.

References

[1] A.H. Guth, Phys. Rev. D 23 (1981) 347.
[2] T. Yanagida, in: O. Sawada, A. Sugamoto (Eds.), Proceedings of the Workshop on Unified Theory and Baryon Number of the Universe, KEK, 1979, p. 95;
M. Gell-Mann, P. Ramond, R. Slansky, in: P. van Niewwenhuizen, D. Freedman (Eds.), Supergravity, North-Holland, Amsterdam, 1979.
[3] A. Hallin, Talk given at XXth International Conference on Neutrino Physics and Astrophysics, Munich, 2002;
M. Smy, Talk given at XXth International Conference on Neutrino Physics and Astrophysics, Munich, 2002;
M. Shiozawa, Talk given at XXth International Conference on Neutrino Physics and Astrophysics, Munich, 2002. For the conference webpage, see http://neutrino2002.ph.tum.de/.
[4] N. Arkani-Hamed, L.J. Hall, H. Murayama, D.R. Smith, N. Weiner, Phys. Rev. D 64 (2001) 115011.
[5] F. Borzumati, Y. Nomura, Phys. Rev. D 64 (2001) 053005.
[6] K. Dick, M. Lindner, M. Ratz, D. Wright, Phys. Rev. Lett. 84 (2000) 4039.
[7] H. Murayama, A. Pierce, hep-ph/0206177.
[8] H. Murayama, T. Yanagida, Phys. Lett. B 322 (1994) 349.
[9] K. Hamaguchi, H. Murayama, T. Yanagida, Phys. Rev. D 65 (2002) 043512.
[10] V.A. Kuzmin, V.A. Rubakov, M.E. Shaposhnikov, Phys. Lett. B 155 (1985) 36.
[11] M. Fukugita, T. Yanagida, Phys. Lett. B 174 (1986) 45.
[12] I. Affleck, M. Dine, Nucl. Phys. B 249 (1985) 361.
[13] H. Murayama, H. Suzuki, T. Yanagida, J. Yokoyama, Phys. Rev. Lett. 70 (1993) 1912.
[14] L. Kofman, A. Linde, A.A. Starobinsky, Phys. Rev. D 56 (1997) 3258.
[15] J.R. Primack, astro-ph/0205391.
[16] M. Kawasaki, T. Moroi, Prog. Theor. Phys. 93 (1995) 879;
M. Kawasaki, K. Kohri, T. Moroi, Phys. Rev. D 63 (2001) 103502.
[17] C. Bennett, et al., Astrophys. J. Lett. 464 (1996) L1.
[18] C.B. Netterfield, et al., astro-ph/0104460.
[19] A.T. Lee, et al., astro-ph/0104459.
[20] N.W. Halverson, et al., astro-ph/0104489.
[21] MAP webpage, http://map.gsfc.nasa.gov.
[22] T. Moroi, T. Takahashi, Phys. Lett. B 522 (2001) 215.
[23] T. Moroi, T. Takahashi, Phys. Rev. D 66 (2002) 063501.
[24] E. Haylo, Phys. Lett. B 387 (1996) 43;
P. Binetruy, G. Dvali, Phys. Lett. B 388 (1996) 241.
[25] W. Hu, astro-ph/9508126.
[26] K. Enqvist, M.S. Sloth, Nucl. Phys. B 626 (2002) 395.
[27] D.H. Lyth, D. Wands, Phys. Lett. B 524 (2002) 5.
[28] J.M. Bardeen, P.J. Steinhardt, M.S. Turner, Phys. Rev. D 28 (1983) 629.
[29] J.R. Bond, A.H. Jaffe, L. Knox, Astrophys. J. 533 (2000) 19.
[30] L. Knox, Phys. Rev. D 52 (1995) 4307.
[31] M. Zaldarriaga, U. Seljak, Astrophys. J. 469 (1996) 437.
[32] RADPACK webpage, http://flight.uchicago.edu/knox/radpack.html.